The nuclear effects in $g_{1^3\mathrm{He}}$ and the Bjorken sum rule for A=3

L.Frankfurt[*], V. Guzey[**], M. Strikman[**]

[*]*School of Physics And Astronomy, Raymond and Beverly Sackler Faculty of Exact Science*

Tel Aviv University, Ramat Aviv 69978, Israel

[**]*Department of Physics, The Pennsylvania State University, University Park, PA 16802*

Abstract

The Bjorken sum rules for the $A = 3$ and $A = 1$ are used as a guide to estimate nuclear effects in extracting $g_{1n}(x, Q^2)$ from the $\vec{e}^3\vec{\mathrm{He}}$ data. We estimate that the combination of the spin depolarization, the nonnucleonic degrees of freedom in the $A = 3$ system and nuclear shadowing is likely to reduce $g_{1^3\mathrm{He}}(x \leq 0.05)$ by $\sim 15\%$ while a significant enhancement of the structure functions $g_{1^3\mathrm{He}}, g_{1^3\mathrm{H}}$ at $x \sim 0.1$ is predicted.

PACS number(s): 25.30.-c, 24.70.+s, 25.10.+s, 29.35.pg

Over the last decade a series of experiments has been performed aiming at measuring the polarized structure functions of protons and neutrons, cf. [1]. The primary motivation was to check the Bjorken sum rule. Recently it has been emphasized that high precision measurements of this sum rule may allow for the accurate determination of $\alpha_s(Q^2)$ [2]. The measurement of $g_{1n}(x)$ involves necessarily nuclear targets. Several experiments using polarized ^{2}H and ^{3}He targets have been performed and several more are in progress. The advantage of the ^{3}He target over the ^{2}H target is that in the first approximation only the neutron is polarized, so that the contribution of the much larger proton structure function $g_{1p}(x, Q^2)$ is small [3].

High precision nonrelativistic calculations of the ^{3}He wave function using realistic nuclear potentials are now available, cf. [4]. They have been applied to analyze the polarized e–^{3}He scattering using the convolution models, where nonnucleonic degrees of freedom in nuclei and nuclear shadowing are neglected [5,6].

The general conclusion is that, similar to the ^{2}H case [7], the major effect of nuclear structure for $x \leq 0.5$ is the depolarization of nucleons in nuclei due to the presence of the higher partial waves. Fermi motion effects do not produce any noticeable x dependence up to $x \sim 0.5$ [5,6]. To avoid dealing with small corrections due to the $\sim 2\%$ polarization of protons in ^{3}He it is convenient to consider the nonsinglet polarized structure functions:

$$g_{1N}^{n.s.}(x, Q^2) \equiv g_{1p}(x, Q^2) - g_{1n}(x, Q^2),$$

and

$$g_{1,A=3}^{n.s.}(x, Q^2) \equiv g_{1^3\mathrm{H}}(x, Q^2) - g_{1^3\mathrm{He}}(x, Q^2).$$

One finds [5,6]

$$g_{1,A=3}^{n.s.}(x, Q^2) = (P_S - \frac{1}{3}P_{S'} + \frac{1}{3}P_D)g_{1N}^{n.s.}(x, Q^2), \tag{1}$$

for $x \leq 0.5$. Here P_S, $P_{S'}$ and P_D are the probabilities of the corresponding components of the neutron wave function in ^{3}He.

For the ratio of the Bjorken sum rule for $A = 3$ to $A = 1$ within the discussed above

impulse approximation the corrections which are proportional to $\alpha_s^n(Q^2)$ cancel out and one obtains:

$$R = \frac{\int_0^1 [g_1^{^3He}(x, Q^2) - g_1^{^3H}(x, Q^2)]dx}{\int_0^1 [g_1^n(x, Q^2) - g_1^p(x, Q^2)]dx} = \frac{G_A(^3H)}{G_A(n)}, \tag{2}$$

independent of Q^2, where we have ignored the higher twist effects. G_A is the axial coupling constant for β decay of the nucleus A. Comparing eqs. (1) and (2) we find

$$R = P_S - \frac{1}{3}P_{S\prime} + \frac{1}{3}P_D. \tag{3}$$

This is perfectly consistent with the expression for $G_A(^3H)$ derived by Blatt back in 1952 [8]:

$$G_A(^3H) = (P_S - \frac{1}{3}P_{S\prime} + \frac{1}{3}P_D)G_A(n). \tag{4}$$

The problem however is that relation (4) is known to be violated experimentally rather significantly. Indeed, realistic 3-nucleon models of ^{3}He and ^{3}H give [4]:

$$P_S - \frac{1}{3}P_{S\prime} + \frac{1}{3}P_D = 1 - (0.0785 \pm 0.0060). \tag{5}$$

Combining the most recent experimental data on $G_A(^3H)/G_V(^3H)$ for tritium β-decay [9] (the data is in good agreement with the previous data [10]) with the value of $G_A(n)/G_V(n)$ from [11]) we obtain

$$\frac{G_A(^3H)}{G_A(n)} = 1 - (0.0366 \pm 0.0030). \tag{6}$$

Hence we conclude that *the use of the convolution model, combined with the 3-nucleon description of $A = 3$ nucleon system, leads to a $\sim 4\%$ violation of the Bjorken sum rule for the scattering of the $A = 3$ systems.* This is consistent with the general expectation that noticeable nonnucleonic degrees of freedom should be present in the $A = 3$ systems.

Nuclear effects for the Bjorken sum rule were first discussed by Close et al [12] and by Kaptari and Umnikov [13]. In particular it was pointed out in Ref. [13] that convolution models and three nucleon description of $A = 3$ system lead to results for g_1 $_{A=3}$

inconsistent with the Bjorken sum rule, though they did not notice consistency of eqs. (3) and (4) which is of importance for our subsequent analysis. This observation was left unnoticed in Refs. [5,6] and in all analyses of the experimental data.

The importance of the Δ-isobar and meson exchange currents for a quantitative explanation of the value of $G_A(^3\text{H})/G_A(n)$ is discussed in literature for a long time, see e.g. Ref. [14]. The recent theoretical analyses of $G_A(^3\text{H})$ [15,16] confirm the conclusion of Ref. [14] that the dominant contribution originates from the admixture of Δ-isobars in ^{3}He and in ^{3}H. They lead to a value of $G_A(^3\text{H})/G_A(n)$ consistent with eq. (6). For example, Ref. [15] gives $G_A(^3\text{H})/G_A(n) = 1 - 0.0378 \pm 0.002$. This implies that the major correction to the impulse approximation calculation of $G_A(A = 3)$ is due to $\Delta \to N$ transitions. Thus a natural mechanism for resolving the discrepancy between the Bjorken sum rule for $A = 3$ and for $A = 1$ targets which is present in the impulse approximation, is the necessity to account for the nondiagonal transitions $\gamma^* N \to \gamma^* \Delta$. No theoretical investigations of this structure function have been done as yet. For the simple case of $g_{1n}^{n.s}$ one can expect the same low x behavior for this structure function as for the diagonal transitions since Regge trajectories with rather close value of intercept couple in this case. Based on $SU(6)$ symmetry, for average $x \sim 0.2 \div 0.3$ we can expect a behavior similar to the diagonal nonsinglet matrix elements. Consequently, we can estimate that the contribution of the $\gamma^* N \to \gamma^* \Delta$ transition to $g_{1,A=3}^{n.s.}$ leads to a change in the ratio $\frac{g_{1,A=3}^{n.s.}(x,Q^2)}{g_{1N}^{n.s.}(x,Q^2)}$ for $x \leq 0.5$ from 1-(0.0785 ± 0.0060) to $G_A(^3\text{H})/G_A(n) = 1 - (0.0366 \pm 0.0030)$. Moreover, treating the Δ-admixture as a perturbation we observe that main contribution to g_1 should originate in the ^{3}He case from the $n \to \Delta^0$ nondiagonal transition and in the ^{3}H case from the $p \to \Delta^+$ nondiagonal transition. In the $SU(6)$ limit, which seems reasonable at least for the valence quark contribution [1]

---

[1] Note that the SMC semiinclusive data [17] seem to indicate that contribution of the sea to g_1 is small down to $x \sim 0.01$.

$$\frac{g_{1\ n\to\Delta^0}(x,Q^2)}{g_{1n}(x,Q^2)} = \frac{g_{1\ p\to\Delta^0}(x,Q^2)}{g_{1p}(x,Q^2)}. \tag{7}$$

Hence up to a small correction due to the contribution of $g_{1p}(x,Q^2)$, the combined effect of nucleon depolarization and nondiagonal contributions is approximately the same for $\frac{g_{1,^3He}(x,Q^2)}{g_{1n}(x,Q^2)}$ and for $\frac{g_{1,A=3}^{n.s.}(x,Q^2)}{g_{1N}^{n.s.}(x,Q^2)}$. We can write in this approximation

$$g_{1^3He}(x,Q^2) = \frac{G_A(^3H)}{G_A(n)}g_{1n}(x,Q^2) + 2p_p(g_{1p}(x,Q^2) + g_{1n}(x,Q^2)), \tag{8}$$

$$g_{1^3H}(x,Q^2) = \frac{G_A(^3H)}{G_A(n)}g_{1p}(x,Q^2) + 2p_p(g_{1p}(x,Q^2) + g_{1n}(x,Q^2)), \tag{9}$$

where $p_p \approx -2.8\%$ is polarization of a proton in ^{3}He. We neglect here contribution of $\Delta^+ \to p$ nondiagonal terms since they effectively merely result in the renormalization of p_p by a factor $\sim \frac{G_A(^3H)}{G_A(n)}$. Experimentally, for small $x \le 0.1$

$$\left|g_{1p}(x,Q^2) + g_{1n}(x,Q^2)\right| \ll \left|g_{1n}(x,Q^2)\right| \left|g_{1n}(x,Q^2)\right|$$

and hence the last term for these x is a very small correction.

Let us calculate nuclear effects specific for the small x physics. At small x, when the coherence length $l = \frac{1}{2m_N x}$ far exceeds the nucleus radius, the virtual photon converts to a quark-gluon configuration h well before the target. In the case of nucleon targets this leads to diffraction in deep inelastic scattering which has recently been observed at HERA. For the nuclear targets this leads to the shadowing phenomenon, for review see [18]. Currently nuclear shadowing in the leading twist is observed experimentally for the sea quark distribution, for the recent review see [19]. There is indirect evidence for the presence of this phenomenon for valence quarks [20]. The presence of gluon shadowing was recently reported based on the analysis of the scaling violation of the $F_{2Sn}(x,Q^2)/F_{2C}(x,Q^2)$ ratio [21].

The phenomenon of shadowing reflects the presence of quark-gluon configurations in γ^* which can interact with cross sections comparable to that of hadrons. A quantitative description of nuclear shadowing phenomenon in deep inelastic scattering was developed

in the color screening models [18,20,22–24], where γ^* converts to a quark-gluon state h which interacts with the nuclear target via multiple color singlet exchanges. The effect of shadowing is determined in these models by the average value of the ratio $\sigma_{eff} = \frac{\langle\sigma^2\rangle}{\langle\sigma\rangle}$, where averaging is taken over different strengths of interaction, that is, over different quark-gluon configurations involved in the transition $\gamma^* \to$ "$hadron\ state$". Numerical analyses of nuclear shadowing for $A \geq 12$ give $\sigma_{eff} \sim 17\ mb$. Similar number follows from the estimate based on the generalization of the optical theorem, cf. [25] to the diffractive processes

$$\sigma_{eff} = \frac{16\pi \frac{d\sigma(\gamma^*+p\to X+p)}{dt}|_{t=0}}{\sigma_{tot}(\gamma^* + p)}. \tag{10}$$

As soon as this parameter is fixed all color singlet models give very similar results for $x \ll \frac{1}{4m_N R_A}$, for a recent discussion and refs. see [26]. We will use this model in the following analysis.

It follows from the formulae of the Glauber approximation that for the case of cross sections which constitute a small fraction of the total cross section, the shadowing effects should be larger. Several examples include shadowing in the parity violating $\vec{p}A$ scattering [27] and shadowing for valence quarks [18]. The underlying physics is quite simple. Let us consider scattering off a heavy nucleus in which one nucleon is polarized. If this nucleon is at a small impact parameter the optical density is high and the cross section of the interaction is not sensitive to its polarization. Hence the cross sections for two polarizations would differ due to large impact parameters only, and therefore shadowing is larger in this case than in the case of the total cross section. Consequently we expect an enhancement of the contribution due to the nuclear shadowing effect to $g_{1A=3}^{n.s}$ as compared to $F_{2\,^3\mathrm{He}}$.

To calculate shadowing for the case of $\vec{e}^{\,3}\vec{\mathrm{He}}$ scattering for $Q^2 \sim Q_0^2 \sim\ few\ GeV^2$ we can consider the difference in the cross sections for the scattering of γ^* with a given helicity (we will not write it explicitly) off $^3\vec{\mathrm{He}}$ with helicities $\pm 1/2$ which we will denote $\pm$. (For larger Q^2 the scaling violation for $F_{2A}(x, Q^2)$, $g_{1A}(x, Q^2)$ can be accounted for using QCD evolution equation.) The cross section can be written in a symbolic form as

$$\sigma_{\gamma^{*3}\vec{H}e_{\pm}} = \sum_h \left| \langle \gamma^* | h \rangle \right|^2 \sigma(h^3 \vec{H}e_{\pm}) \tag{11}$$

We substituted the integral over the hadronic state by its value at an average point that has an interaction with a nucleon σ_{eff}, and mass of the state h is $M^2 = Q^2$ [18].

For simplicity we consider the model where all nucleons in the nucleus of ^{3}He are in the $S-$state and hence only the neutron is polarized. However we expect that nuclear shadowing effects should lead to a universal factor weakly dependent on the form of the wave function of the nucleus. To calculate $\sigma(h^3 \vec{He}_{\pm})$ we use the modified Glauber method [28] which takes into account the fact that the longitudinal momentum transferred in the transition $\gamma \to h$ is $q_{\|} = \frac{Q^2 + M_h^2}{2q_0}$. Within the above approximation we have $q_{\|} = 2m_N x$. If we include all possible permutations of the nucleons, we can write the modified profile function in the following form

$$\Gamma(\vec{\rho}, \vec{r_{1t}}, \vec{r_{2t}}, \vec{r_{3t}}) = \Gamma_n(\vec{\rho} - \vec{r_{1t}}) + 2\,\Gamma_p(\vec{\rho} - \vec{r_{1t}}) - 4\,\Gamma_n(\vec{\rho} - \vec{r_{1t}})\Gamma_p(\vec{\rho} - \vec{r_{2t}})\Theta(z_2 - z_1)e^{iq_{\|}(z_1 - z_2)}$$
$$- 2\,\Gamma_p(\vec{\rho} - \vec{r_{1t}})\Gamma_p(\vec{\rho} - \vec{r_{2t}})\Theta(z_2 - z_1)e^{iq_{\|}(z_1 - z_2)}$$
$$+ 6\,\Gamma_n(\vec{\rho} - \vec{r_{1t}})\Gamma_p(\vec{\rho} - \vec{r_{2t}})\Gamma_p(\vec{\rho} - \vec{r_{3t}})\Theta(z_2 - z_1)\Theta(z_3 - z_2)e^{iq_{\|}(z_1 - z_3)}. \tag{12}$$

In these estimates we have accounted only for elastic rescatterings of the state $|h\rangle$. It is a reasonable approximation at moderate Q^2. The scattering amplitudes f are related to $\Gamma(\vec{\rho})$ as

$$f_{hp(n)}(q) = \frac{ik}{2\pi} \int e^{-i\vec{\rho}\cdot\vec{q_t}} \Gamma_{p(n)}(\vec{\rho}) d^2\vec{\rho} \tag{13}$$

The ^{3}He wave function is taken in a simple form (S-state), which works well in the Glauber calculations of elastic $p\,^4$He scattering [29]: $|\Psi|^2 \propto \prod_{l=1}^{l=3} \exp(-\vec{r}_l^2/2\alpha)\delta^3(\sum \vec{r}_l)$. So, only the neutron is polarized in this approximation. The numerical value of the slope was fixed to reproduce the e.m. form factor of ^{3}He: α=27 GeV^{-2}. Within the described above approximation the t dependence of the amplitude $hN \to hN$ is the same as for the amplitude $\gamma^* + N \to h + N$. Hence on the basis of current experience, we write

$$f_{hp}(q_t) = is\sigma_p e^{-\beta/2\,q_t^2}(1 + i\eta) \tag{14}$$

7

$$f_{hn}^{\pm}(q_t) = is\sigma_n^{\pm}e^{-\beta/2\,q_t^2}(1 + i\eta_{\pm}),\tag{15}$$

where $\eta = Re\,f_{hp}/Im\,f_{hp}$, $\eta_{\pm} = Re\,f_{hn}^{\pm}/Im\,f_{hn}^{\pm}$, $\beta \approx 6$ GeV^{-2}. Note that since we are concerned here with the x and Q^2 ranges corresponding to the energies relevant to the current measurements of g_{1n}, in estimating β from the HERA data we take into account a weak energy dependence in the slope expected for the Regge pole approximation. We also assume that the slope for the spin dependent amplitude is the same as for the spin independent amplitude. Since both slopes are much smaller than the nuclear form factor slope our result is not sensitive to the value of β. Finally we obtain for the total cross section

$$\sigma_T^{\pm} = \sigma_n^{\pm} + 2\sigma_p - \frac{\sigma_p^2 e^{-\alpha q_{\parallel}^2}}{8\pi(\alpha + \beta)}\left(1 - \eta^2 - 2\eta\sqrt{\frac{4\alpha}{\pi}}\cdot q_{\parallel}\right)$$
$$- \frac{\sigma_n^{\pm}\sigma_p e^{-\alpha q_{\parallel}^2}}{4\pi(\alpha + \beta)}\left(1 - \eta_{\pm}\eta - (\eta_{\pm} + \eta)\sqrt{\frac{4\alpha}{\pi}}\cdot q_{\parallel}\right) + \frac{1}{48\pi^2(\alpha + \beta)^2}\sigma_p^2\sigma_n^{\pm}e^{-\alpha q_{\parallel}^2}.\tag{16}$$

In the third term, which is numerically small, we neglected the corrections due to the real part of the amplitude and higher order corrections in $q_{\parallel}$. Using eq.(16) we evaluate the shadowing in the case of the unpolarized target:

$$\frac{F_{2A=3}(x, Q_0^2)}{3F_{2N}(x, Q_0^2)} = 1 - \frac{\sigma_{eff}}{8\pi(\alpha + \beta)}exp(-\alpha q_{\parallel}^2)\left(1 - \eta^2 - 2\eta\sqrt{\frac{4\alpha}{\pi}}\cdot q_{\parallel}\right)$$
$$+ \frac{\sigma_{eff}^2}{144\pi^2(\alpha + \beta)^2}\exp(-\alpha q_{\parallel}^2).\tag{17}$$

For $g_{1^3He}(x, Q^2)$ we obtain:

$$\frac{g_{1^3He}(x, Q_0^2)}{g_{1n}(x, Q_0^2)} = \frac{\sigma_T^{+}(e^3He) - \sigma_T^{-}(e^3He)}{\sigma_T^{+}(en) - \sigma_T^{-}(en)} =$$
$$1 - \frac{\sigma_p exp(-\alpha q_{\parallel}^2)}{4\pi(\alpha + \beta)}(1 - K) + \frac{\sigma_p^2 exp(-\alpha q_{\parallel}^2)}{48\pi^2(\alpha + \beta)^2}.\tag{18}$$

Here K is given by

$$K = (\sqrt{\frac{4\alpha}{\pi}}q_{\parallel} + \eta)\cdot\frac{\sigma_n^{+}\eta_{+} - \sigma_n^{-}\eta_{-}}{\sigma_n^{+} - \sigma_n^{-}}.\tag{19}$$

Similar expressions are valid for the ratios g_{1^3H}/g_{1p} and $g_{1A=3}^{n.s.}/g_{1N}^{n.s.}$. Factors η, $\eta_{\pm}$ are small because the vacuum exchange dominates in rescattering amplitudes. This

approximation may become dangerous for $x \leq 10^{-3}$ where $F_{2p}(x, Q^2)$ starts to increase fast with decrease of x. At the same time the factor $(\sigma_n^+ \eta_+ - \sigma_n^- \eta_-)/(\sigma_n^+ - \sigma_n^-)$, which is determined by the phase of the secondary Regge trajectories which dominate in $g_{1N}(x, Q^2)$ for small x, could be of order unity. However, its contribution is suppressed for small x by the factor of $q_{\parallel}$ and the small value of η for $x \geq 10^{-3}$. Hence in our estimates we neglect the contributions of the real part of all amplitudes. Uncertainties of this approximation will be analyzed elsewhere.

One can see from the comparison of eqs.(17) and (18) that shadowing for the case of the polarized cross section is larger by approximately a factor of two. This result justifies the above qualitative discussion. Eq.(18) leads to $\frac{g_{1^3He}(x, Q_0^2)}{g_{1n}(x, Q_0^2)} \approx 0.9$ for $x \leq 0.03$. Obviously, nuclear shadowing changes the contribution at small x to the Bjorken sum rule. As in the case for valence quarks (baryon sum rule) and gluon distributions (momentum sum rule) the compensating positive contribution to $g_{1A=3}^{n.s.}(x, Q^2)$ related to the projectile interaction with two nucleons should be located at $x \sim 0.1$, cf. discussion in [18]. Hence we model this enhancement by requiring that (i) the positive contribution to $g_{1A=3}^{n.s.}(x, Q^2)$ compensates the contribution due to shadowing in $\int_0^1 g_{1A=3}^{n.s.}(x, Q^2)dx$, (ii) does not affect the region where shadowing is saturated ($x \leq 0.03$), (iii) it is concentrated for $x \leq 0.15$. An example of this fit is given in Fig.1 by a dashed line. One can see that typically the resulting enhancement is of the order $10 \div 15\%$.

Thus we conclude that there are two new effects modifying the picture of nuclear effects for $g_{1^3He}(x, Q^2)$ based on the nonrelativistic model of the nucleus: the nonnucleonic degrees of freedom and nuclear shadowing. Based on the additive quark model one expects that in the high-energy limit $\sigma_{tot}(|h\rangle \Delta) \approx \sigma_{tot}(|h\rangle N)$. Hence shadowing effect should be approximately the same for the nondiagonal contribution. Difference in the shadowing for higher partial waves maybe somewhat larger due to smaller radius of these components of the wave function. However since the shadowing effect is rather small this effect does not exceed the product of the shadowing correction and depolarization effect. Thus it leads to an uncertainty in the discussed ratio of less than 1%. Hence

9

the two discussed effects in the first approximation contribute multiplicatively to the modification of $g_{1A=3}^{n.s.}$ which is given by the solid curve. Our expectations for the g_{1^3He}/g_{1n} and g_{1^3H}/g_{1p} ratios are practically the same. It is noticeably different from the $\sim 8\%$ depolarization effect obtained in the model [5,6] (the dashed-dotted line), where these effects were neglected. Substantial model dependence of the nuclear effects introduces significant uncertainties in the extraction of g_{1n} from the ^{3}He data, especially for $x \leq 0.2$. The detailed procedure of extraction of g_{1n} would involve separate modeling the Δ-contribution to $g_{1^3\mathrm{He}}$ and $g_{1^3\mathrm{H}}$ going beyond the $SU(6)$ -symmetry approximation and calculation of shadowing effects including effects of higher partial waves. We will consider these effects elsewhere.

In this x-range ^{2}H targets may have certain advantages since in this case nonnucleonic admixtures are much smaller due to weaker binding and zero isospin. The shadowing effects are also smaller for $g_{1^2\mathrm{H}}$ by a factor of $\sim 0.4 \div 0.5$ [30]. Besides, in the first approximation $g_{1p}(x, Q^2) \approx -g_{1n}(x, Q^2)$ for small x in which case shadowing does not affect the extraction of g_{1n}.

Further studies are necessary to work out the x-dependence of the contribution of the nonnucleonic degrees of freedom and to develop a dynamic mechanism of the enhancement effects for $g_{1A=3}^{n.s.}$.

Also it would be interesting to check the predicted patterns for the screening-enhancement in independent experiments with other polarized nuclei where the polarization is carried predominantly by a proton. Obviously, the heavier the nucleus, the larger the effect. Another interesting question is the relation of shadowing to quenching of G_A/G_V for heavy nuclei. If there were no enhancement at moderate x associated with shadowing at small x, it would lead to renormalization of G_A/G_V by 10-20 %.

REFERENCES

[1] J.Ellis and M.Karliner, preprint CERN-TH/95-279.

[2] J.Ellis and M.Karliner, Phys.Lett. **B341**, 397 (1995).

[3] R.Blankleider and R.M.Woloshin, Phys.Rev. **C29** 538 (1984); R.M.Woloshin, Nucl.Phys. **A495** 749 (1985).

[4] J.L.Friar, B.F.Gibson, G.L.Payne, A.M.Bernstein, T.E.Chupp, Phys. Rev. **C42** ,2310 (1990); J.Friar private communication 1993.

[5] C. Ciofi degli Atti, S. Scopetta, E. Pace, and G. Salme, Phys.Rev. **C48**, 968 (1993).

[6] R.W.Schultze and P.U.Sauer, Phys.Rev. **C48** 38 (1993).

[7] L.Frankfurt and M.Strikman, Nucl.Phys. **A405**, 557 (1983).

[8] J.M.Blatt, Phys.Rev. **89**, 86 (1952); L.M.Delves and A.C.Phillips, Rev.Mod.Phys. **41**, 497 (1969).

[9] B.Budick, J.Chen and H.Lin, Phys.Rev.Lett. **67** ,2630 (1991).

[10] J.J.Simpson, Phys.Rev **C35**, 752 (1987)

[11] The Rev. of Particle Properties, Phys. Rev **D50**, 1173 (1994)

[12] F.E. Close, R.G. Roberts, G.G. Ross, Nucl.Phys. **B296**,582 (1988).

[13] L.P. Kaptari, A.Yu. Umnikov, Phys.Lett. **B240** 203 (1990).

[14] T.Ericson and W.Weise, Pions and nuclei, Clarendon Press, 1988.

[15] T.Y.Saito, Y.Wu,S.Ishikawa, and T.Sasakawa, Phys.Lett. **B242**, 12 (1990).

[16] J.Adam Jr, Ch.Hajduk, H.Henning, P.U.Sauer, and E.Truhlik, Nucl.Phys. **A531**, 623 (1991).

[17] SMC, B. Adeva et al., Phys.Lett.B369,93(1996).

[18] L. L. Frankfurt and M. Strikman, Phys. Rep. **160**, 235 (1988) and Nucl. Phys. **B316**, 340 (1989).

[19] M.Arneodo, Physics Reports **240**, 301 (1994).

[20] L. L. Frankfurt, M. Strikman and S. Liuti, Phys. Rev. Lett. **65**, 1725 (1990)

[21] T.Gousset and H.J.Pirner, preprint hep-ph/9601242

[22] S.J.Brodsky and H.J.Lu, Phys.Rev.Lett. **64**, 1342 (1990).

[23] N. N. Nikolaev and B. G. Zakharov, Z. Phys. **C49**, 607 (1991).

[24] G.Piller, W.Ratzka, and W.Weise, Z.Phys. **A352**, 427 (1995).

[25] H. Miettinen and J. Pumplin, Phys. Rev. **D18**, 1696 (1978); Phys. Rev. Lett. **42**, 204 (1979).

[26] L. Frankfurt and M. Strikman, Proc. of Workshop on Deep Inelastic Scattering and QCD, Paris, France, 1995, Ed. J.F.Laporte and Y.Sirois, p.485.

[27] L.L.Frankfurt and M.I.Strikman, Phys.Lett. **107B**, 99 (1981); Progr. Part.Nucl.Phys, **27**, 135 (1991).

[28] T.H.Bauer, R.D.Spital, D.R.Yennie and F.M.Pipkin, Rev. Mod.Phys. **50**, 261 (1978).

[29] E.M.Levin and M.I.Strikman, Sov.J.Nucl.Phys. **23**, 216 (1976).

[30] A. Berera, L. Frankfurt and M. Strikman, in preparation

[31] L.Frankfurt, V.Guzey and M.Strikman, hep-ph/9602301.

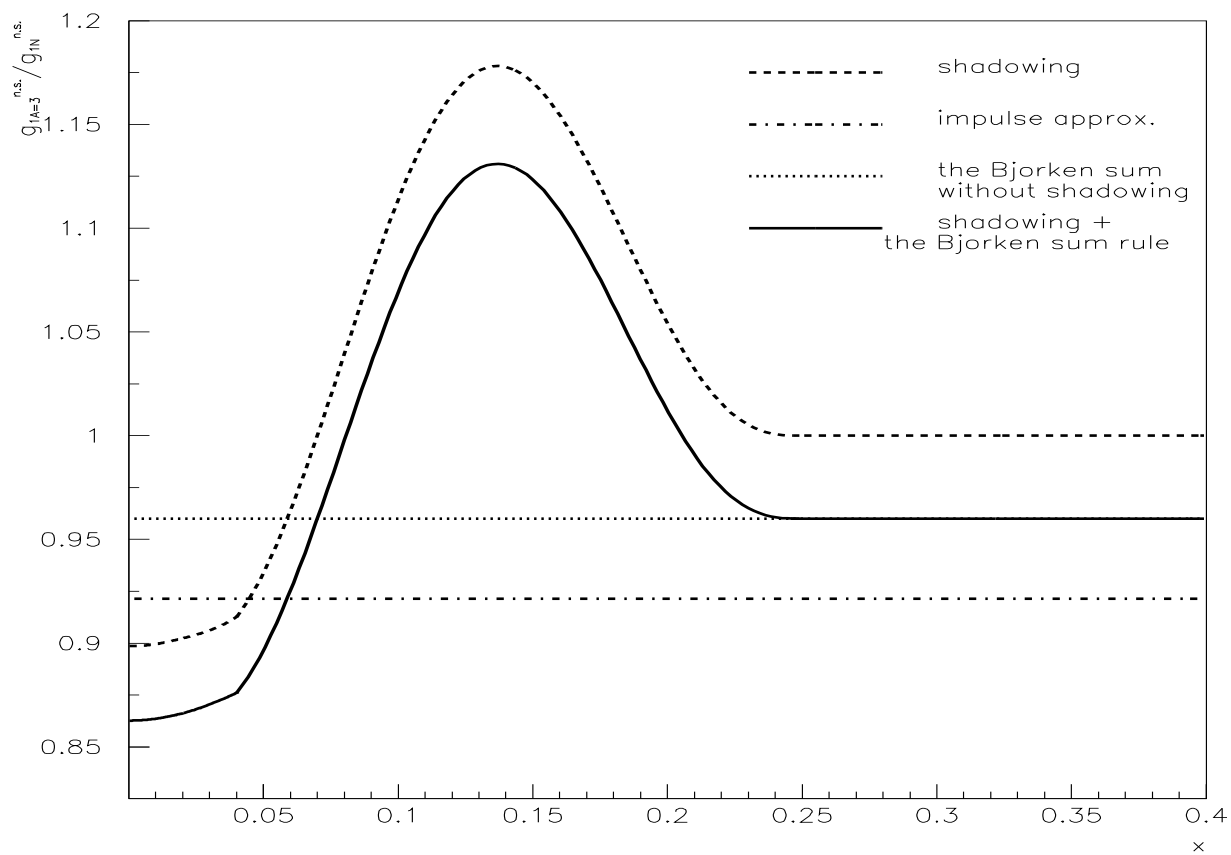FIG. 1. $g_{1A=3}^{n.s.}/g_{1N}^{n.s.}$ as a function of x. The dashed line represents nuclear shadowing at small x. The solid line is the result of the fit constrained to preserve the Bjorken sum rule.